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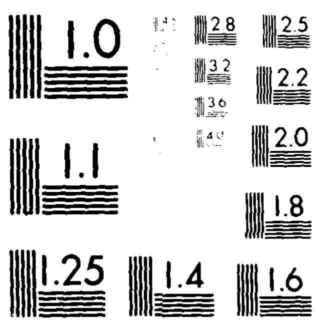

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On the Spatial Power Spectrum of the $E \times B$ Gradient Drift Instability in Ionospheric Plasma Clouds.

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Geophysical and Plasma Dynamics Branch
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It is shown that the two dimensional power law spatial power spectra of magnetic field line integrated Pedersen conductivity perturbations in the plane perpendicular to the magnetic field recently computed from numerical simulation studies of ionospheric plasma coluds are consistent with conservation laws implied by the fundamental plasma fluid equations describing the $\underline{E} \times \underline{B}$ gradient drift instability in these ionospheric plasma clouds.

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ON THE SPATIAL POWER SPECTRUM OF THE $\underline{E} \times \underline{B}$ GRADIENT DRIFT INSTABILITY IN IONOSPHERIC PLASMA CLOUDS

Experimental studies of plasma clouds in the ionosphere [Rosenberg, 1971; Davis et al., 1974; Baker and Ulwick, 1978] have yielded much data describing ambient ionospheric conditions, e.g., electric and magnetic fields. The characteristic initial steepening, elongation, and striation of $\underline{E} \times \underline{B}$ drifting plasma clouds have been studied by applying the linear theory of the $\underline{E} \times \underline{B}$ gradient drift instabiltiy, originally developed for laboratory gas discharges [Simon, 1963] to plasma clouds geometries [Haerendel et al., 1967; Linson and Workman, 1970; Völk and Haerendel, 1971; Perkins et al., 1973]. More recently Chaturvedi and Ossakow [1979] has presented arguments based on the nonlinear two-dimensional coherent mode coupling of two Fourier modes of plasma cloud density, following the work of Rognlien and Weinstock [1975], to explain the nonlinear stabilization and resultant saturated amplitudes and power spectra of the long wavelength $\underline{E} \times \underline{B}$ gradient drift instability in ionospheric plasma clouds.

Numerical simulations [Zabusky et al., 1973; Lloyd and Haerendel, 1973; Goldman et al., 1974; Doles et al., 1976; Ossakow et al., 1975, 1977] of the interaction of barium plasma clouds with the ionosphere have reproduced not only many of the gross observational features of plasma cloud evolution, but also their spatial power spectra [Scannapieco et al., 1976], minimum striation scale sizes [McDonald et al., 1978, 1980] and outer scale size or correlation length [Keskinen et al., 1980a; Keskinen and Ossakow, 1980]. In addition numerical

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simulation studies of the E x B gradient drift instability in local unstable regions of ionospheric plasma clouds [Keskinen et al., 1980b] have yielded spatial power spectra and saturated wave amplitudes consistent with experimental values.

However, to date, an analytical description of the two-dimensional spatial power spectra of the E x B gradient drift instability in ionospheric plasma clouds over a wide range of wavelengths has not been discussed in detail. In this paper we show that the spatial power spectra of ionospheric plasma clouds recently computed from numerical simulation studies and experimental observations are consistent with conservation laws implied by the fundamental fluid equations modeling the E x B gradient drift instability in ionospheric plasma clouds.

For wavelengths much greater than the ion gyroradius (approximately 10 meters for Ba^+ ions in the twilight F region), the dynamics of the plasma cloud and background ionosphere can be studied in the fluid approximation [Völk and Haerendel, 1971; Perkins et al., 1973; Zabusky et al., 1973; Ossakow et al., 1975]. For large clouds (large magnetic field line integrated Pedersen conductivity compared with that of the background ionosphere), the cloud interaction with the background ionosphere (second level) can be neglected [Haerendel et al., 1967]. Furthermore, due to the very high conductivity along the magnetic field lines (typically $\sigma_p/\sigma_{\parallel} \approx 10^{-6}$ with σ_p and $\sigma_{\parallel}$ the Pedersen and parallel conductivities respectively) the field lines may be regarded as equipotentials and, as a result, a

field line integrated (two-dimensional) model is justified [Volk and Haerendel, 1971; Perkins et al., 1973]. In reality, an artificially injected plasma cloud will, initially, be two-dimensional in the plane perpendicular to the magnetic field. But a two-dimensional plasma cloud of finite extent is not stationary [Dungey, 1958] and will distort as it convects. However, for striation scale sizes less than the ambient cloud Pedersen conductivity gradient scale length and on the faster time scale on which striations develop, the assumption of one-dimensionality is justified. As a result, in what follows we confine our attention to local regions of initially slablike (one-dimensional) ionospheric plasma clouds. This will also facilitate comparision with numerical simulation studies.

By adopting a Cartesian coordinate system (x,y,z) with magnetic field $B\hat{\underline{z}}$, ambient electric field $E_o\hat{\underline{y}}$, and after transforming to a frame drifting with velocity $\underline{V}_o = (cE_o/B)\hat{\underline{x}}$, the two-dimensional model equations for the magnetic field line integrated plasma cloud Pedersen conductivity $\Sigma(x,y)$ and the self-consistent plasma cloud electrostatic potential $\delta\varphi(x,y)$ can be written

$$\frac{\partial \Sigma}{\partial t} + \frac{c}{B}\,\hat{\underline{z}} \times \nabla\delta\varphi \cdot \nabla\Sigma = D\nabla^2\Sigma \tag{1}$$

$$\nabla \cdot (\Sigma\ \nabla\delta\varphi) = \underline{E}_o \cdot \nabla\Sigma \tag{2}$$

where $\nabla\delta\varphi = -\underline{E}(x,y) + \underline{E}_o$, $\underline{E}(x,y)$ the total electrostatic field, c is the speed of light, $\nabla \equiv (\partial/\partial x,\ \partial/\partial y)$ and D is the cross-field diffusion coefficient [Perkins et al., 1973] given by $2(\nu_e/\Omega_e)(ck_BT/eB)$ with T the ion and electron temperature,

ν_e the sum of electron collision frequencies with plasma cloud ions and ambient neutrals, k_B is Boltzmann's constant, and Ω_e the electron gyrofrequency. All other symbols retain their conventional meaning. Equation (1) results from the magnetic field line integration (along the z-direction) of the ion continuity equation, while equation (2) is derived from current conservation $\nabla \cdot \underline{J} = 0$.

By linearizing (1) and (2) and assuming fluctuations in magnetic field line integrated-Pedersen conductivity and cloud potential $\delta\Sigma, \delta\varphi \propto \exp\left[i(k_y y + k_x x) + \gamma_{\underline{k}} t\right]$ with $\underline{k} \cdot \underline{B} = 0$, $kL >> 1$, one finds the usual $\underline{E} \times \underline{B}$ growth rate

$$\gamma_{\underline{k}} = (cE_o/BL)(k_y/k)^2 - Dk^2 \tag{3}$$

where $k^2 = k_x{}^2 + k_y{}^2$, $L^{-1} = \partial \ln \Sigma_o / \partial x$. For $cE_o/B = 100$ m/sec, $L = 6$ km, $D = 1 m^2/sec$, the critical wavelength λ_c $(\gamma_{\underline{k}} = 0) \simeq 60$ m.

We will now show that several features of the spatial power spectra of the $\underline{E} \times \underline{B}$ gradient drift instabiltiy in ionospheric plasma clouds as observed in experiments and recent numerical simulations can be derived by considering the conservation laws implied by (1) and (2). We can write (1) and (2), after making the separation $\Sigma(x,y) = \Sigma_o(x) + \delta\Sigma(x,y)$, as follows:

$$\frac{\partial \delta\Sigma}{\partial t} + \frac{c}{B} \hat{\underline{z}} \times \nabla\delta\varphi \cdot \nabla\Sigma_o + \frac{c}{B} \hat{\underline{z}} \times \nabla\delta\varphi \cdot \nabla\delta\Sigma = D\nabla^2\delta\Sigma \tag{4}$$

$$\Sigma_o \nabla^2 \delta\varphi + \delta\Sigma \nabla^2 \delta\varphi + \nabla\Sigma_o \cdot \nabla\delta\varphi + \nabla\delta\Sigma \cdot \nabla\delta\varphi = \underline{E}_o \cdot \nabla\delta\Sigma \tag{5}$$

We first multiply (4) by $\delta\Sigma$ and integrate over all x and y finding

$$\frac{1}{2}\frac{\partial}{\partial t}\int d^2x \quad (\delta\Sigma)^2 + \frac{c}{B}\int d^2x \quad \delta\Sigma \ \hat{\underline{z}} \ x\nabla\delta\varphi\cdot\nabla\Sigma_o$$

$$+ \frac{c}{B}\int d^2x \ \delta\Sigma \ \hat{\underline{z}} \ x\nabla\delta\varphi\cdot\nabla\delta\Sigma - D\int d^2x \ \delta\Sigma\nabla^2\delta\Sigma = 0 \qquad (6)$$

where $d^2x \equiv dxdy$. The third term on the left hand side vanishes since

$$\int d^2 x \ \delta\Sigma \ \hat{\underline{z}}x\nabla\delta\omega\cdot\nabla\delta\Sigma$$

$$= \frac{1}{2}\int d^2x \ \hat{\underline{z}} \ x \ \nabla\delta\varphi\cdot\nabla(\delta\Sigma)^2$$

$$= \frac{1}{2}\int d^2x \ \nabla\cdot\left[(\delta\Sigma)^2\hat{z}x\nabla\delta\varphi\right]$$

$$= 0 \text{ if } (\delta\Sigma)^2 \rightarrow 0 \text{ as } x,y \rightarrow \infty .$$

since $\nabla\cdot\left[(\delta\Sigma)^2 \ \hat{z} \ x\nabla\delta\varphi\right] = \hat{z} \ x \ \nabla\delta\varphi\cdot\nabla(\delta\Sigma)^2$. After making the following Fourier expansion of $\delta\Sigma$ and $\delta\varphi$

$$\delta\Sigma(x,y) = \int d^2k' \delta\Sigma_{\underline{k}'} \ e^{i\underline{k}'\cdot\underline{x}}$$

$$\delta\varphi(x,y) = \int d^2k \ \delta\varphi_{\underline{k}} \ e^{i\underline{k}\cdot\underline{x}}$$

the second term on the left hand side of (6) can be written

$$\frac{c}{B}\int d^2x\int d^2k' \ \delta\Sigma_{\underline{k}'} \ e^{i\underline{k}'\cdot\underline{x}} \ \hat{\underline{z}}x\nabla\left(\int d^2k \ \delta\varphi_{\underline{k}} \ e^{i\underline{k}\cdot\underline{x}}\right)\cdot\nabla\Sigma_o$$

$$= \frac{c}{B}\iint d^2kd^2k' \ i \ \hat{\underline{z}} \ x \ \underline{k}\cdot\nabla\Sigma_o \ \delta\Sigma_{\underline{k}'} \ \delta\varphi_{\underline{k}}\int d^2x \ e^{i(\underline{k}'+\underline{k})\cdot\underline{x}}$$

$$= \frac{c}{B}\iint d^2k \ d^2k' \ i \ \hat{\underline{z}} \ x \ \underline{k}\cdot\nabla\Sigma_o \ \delta\Sigma_{\underline{k}'} \ \delta\varphi_{\underline{k}}\delta(\underline{k}'+\underline{k})$$

$$= \frac{c}{B}\int d^2k \ i \ \hat{z} \ x \ \underline{k}\cdot\nabla\Sigma_o\delta\Sigma_{-\underline{k}} \ \delta\varphi_{\underline{k}}$$

where $\delta(\underline{x})$ is the two-dimensional delta function. The first and fourth terms in equation (6) are reduced in a similar

manner

$$\frac{1}{2}\ \frac{\partial}{\partial t}\int d^2x(\delta\Sigma)^2 = \frac{1}{2}\ \frac{\partial}{\partial t}\int d^2k\ |\delta\Sigma_{\underline{k}}|^2$$

and

$$D\int d^2\,x\,\delta\Sigma\ \nabla^2\delta\Sigma = -\ D\int d^2k\ \ k^2|\delta\Sigma_{\underline{k}}|^2$$

All that remains is to find an expression for $\delta\varphi_{\underline{k}}$ to evaluate the second term on the left hand side. Since we are considering only quadratic nonlinearities only a linear relation of $\delta\varphi_{\underline{k}}$ to $\delta\Sigma_{\underline{k}}$ is needed and one can use the linearized version of equation (5) to find $\delta\varphi_{\underline{k}}$ in terms of $\delta\Sigma_{\underline{k}}$.

$$\Sigma_o\ \nabla^2\delta\varphi + \nabla\Sigma_o\cdot\nabla\delta\varphi = \underline{E}_o\cdot\nabla\delta\Sigma \tag{5a}$$

Inserting Fourier expansions of $\delta\Sigma$ and $\delta\varphi$ into (5a) we find in a straightforward manner

$$\delta\varphi_{\underline{k}} \cong -\ (i\underline{k}\cdot\underline{E}_o/k^2\Sigma_o)\delta\Sigma_{\underline{k}} \tag{7}$$

Substituting into the Fourier transformed version of the second term on the left hand side of equation (6) gives

$$\frac{c}{B}\int d^2k\ (\hat{z}\ x\ \underline{k}\cdot\nabla\Sigma_o\ \underline{k}\ \cdot\underline{E}_o/k^2\Sigma_o)\quad|\delta\Sigma_{\underline{k}}|^2$$

Equation (6) can than be written as

$$\frac{1}{2}\ \frac{\partial}{\partial t}\int d^2k\ \ P_{\underline{k}} = \int d^2k\ \gamma_{\underline{k}}\ P_{\underline{k}} \tag{8}$$

where $P_{\underline{k}} \equiv |\delta\Sigma_{\underline{k}}|^2 = |\delta\Sigma(k_x,k_y)|^2$ and $\gamma_{\underline{k}}$ is the linear growth rate given by

$$\gamma_{\underline{k}} = \frac{c}{B} \left(\underline{k} \cdot \underline{E}_o \; \underline{\hat{z}} \times \underline{k} \cdot \nabla \Sigma_o / k^2 \Sigma_o\right) - Dk^2$$

$$= (cE_o/BL) \cos^2\Theta - Dk^2$$

where $L^{-1} = (1/\Sigma_o)(\partial\Sigma_o/\partial x)$ and Θ is the angle defined by $\underline{k}$ and $\underline{E}_o$. We note that the nonlinear term does not enter into (8); the nonlinear term thus conserves energy and may be thought of as a bridge between the stable and unstable regions of wavenumber space. This property is also shared by several other nonlinear plasma physics problems, e.g., the two-dimensional electrostatic guiding center plasma [Montgomery, 1976], ion-acoustic shock formation [Ott et al., 1973], and two-dimensional turbulence in the equatorial electrojet [Keskinen et al., 1979; Sudan and Keskinen, 1977, 1979].

In the steady state ($\partial/\partial t = 0$) we can write (8) in polar coordinates ($k^2 = k_x^{\,2} + k_y^{\,2}$, $\tan\Theta = k_x/k_y$) assuming the spatial power spectrum extends from k_{min} to k_{max}

$$\int_{k_{min}}^{k_c} dk\; k \int_o^{2\pi} d\Theta\; \gamma_{\underline{k}}^g\, P_k = \int_{k_c}^{k_{max}} dk\; k \int_o^{2\pi} d\Theta\; \gamma_{\underline{k}}^d\, P_{\underline{k}} \tag{9}$$

with $\gamma_{\underline{k}}^g = (cE_o/BL)\cos^2\Theta$ and $\gamma_{\underline{k}}^d = Dk^2$. As a result the total power generated in the unstable range of wavenumbers between k_{min} and k_c, the critical wavenumber, must be balanced by the power dissipated in the stable regime between k_c and k_{max} under steady state conditions.

Our approach will be to make an ansatz for $P(k,\Theta)$ and show its consistency with (9). Previous numerical simulation studies

[Scannapieco et al., 1976; Keskinen et al., 1980a; Keskinen and Ossakow, 1980] have shown that the one-dimensional transverse integrated spatial power spectra of magnetic field line integrated Pedersen conductivity perturbations in the plane perpendicular to the magnetic field of initially slablike ionospheric plasma clouds can be represented by a power law, i.e.,

$$\int dk_x \, |\delta\Sigma(k_x,k_y)|^2 \propto (k_y^2 + k_{oy}^2)^{-n_y/2}$$

and

$$\int dk_y \, |\delta\Sigma(k_x,k_y)|^2 \propto (k_x^2 + k_{ox}^2)^{-n_x/2}$$

with $n_x \simeq 2$, $n_y \simeq 2\text{-}3$, and $2\pi/k_o$ the outer scale size where the x-axis denotes the direction of plasma cloud $\underline{E} \times \underline{B}$ drift. These power spectral indices are in agreement with those derived from experimental observations [Baker and Ulwick, 1978; Kelley et al., 1979] of ionospheric plasma clouds. Recently, the angular dependence of $P(k,\Theta)$ has been studied [Keskinen and Ossakow, 198C] in the nonlinear steady state regime for initially slablike ionospheric plasma clouds under different sets of initial conditions. In polar coordinates, these studies show that $P(k,\Theta)$ is anisotropic with maximum power along and near the y-direction (perpendicular to $\underline{E}_o \times \underline{B}$ direction) and suggest that

$$P(k, \Theta) \cong P_o \cos^m\Theta \, (1 + (k/k_o)^2)^{-(n+1)/2}$$

with $m>0$, P_o = const, and $2\pi/k_o$ is the outer scale size. We note that this form for $P(k,\Theta)$ maximizes in the linearly most unstable direction ($\Theta = 0$ where Θ is defined by $\underline{k}$ and $\underline{E}_o$). This anisotropy is a reflection of the fact that nonlinear plasma

cloud striations are elongated much more in the $\underline{E}_o \times \underline{B}$ direction and this result is also observed in local simulations of ionospheric plasma clouds [Keskinen et al., 1980b]. The anisotropic nature of the two-dimensional spatial power spectra of ionospheric plasma clouds is also consistent with arguments put forth by Chaturvedi and Ossakow [1979] in discussing the nonlinear stabilization of the $\underline{E} \times \underline{B}$ gradient-drift instability in ionospheric plasma clouds. In addition, this anisotropy is also observed in experimental plasma cloud studies [M.C. Kelley, private communication, 1980].

We substitute this ansatz for $P(k,\theta)$ into the left hand side of (9) obtaining

$$(P_o cE_o/BL)\int_o^{2\pi} d\theta \cos^{m+2}\theta \int_{k_{min}}^{k_c} dk\, k\, (1+(k/k_o)^2)^{-(n+1)/2}$$

$$\simeq (P_o cE_o/BL)\; 2\sqrt{\pi}\; \frac{\Gamma(\frac{m+3}{2})}{\Gamma(\frac{m}{2}+2)} \frac{k_o^2}{1-n}\left[(1+(\frac{k_c}{k_o})^2)^{(1-n)/2} - (1+(\frac{k_{min}}{k_o})^2)^{(1-n)/2} \right]$$

$$\simeq (P_o cE_o/BL)\; 2\sqrt{\pi}\; \frac{\Gamma(\frac{m+3}{2})}{\Gamma(\frac{m}{2}+2)} \frac{k_o^2}{n-1}$$

where $\Gamma(p)$ is the Gamma-function defined by

$$\Gamma(p) = \int_o^\infty dx\; x^{p-1} e^{-x}$$

and we have assumed $k_c > k_o$, $k_{min} < k_o < k_c$ and $n>1$. Similiarly the right hand side of (9) gives

$$P_o D \quad 2\sqrt{\pi} \quad \frac{\Gamma(\frac{m+1}{2})}{\Gamma(\frac{m}{2}+1)} \quad \frac{k_o^{\,4}}{3-n} \quad \left(\frac{k_{max}}{k_o}\right)^{3-n}$$

We note that since the left and right hand sides of (9) must be positive, the spectral index $1<n<3$. This implies that the steady state spatial power spectrum can be represented by a power law. Equating the power generated in the unstable region to the power dissipated and solving for the outer scale wavenumber k_o we find

$$\left(\frac{k_o}{k_{max}}\right)^{n-1} = \frac{3-n}{n-1} \quad \frac{\Gamma(\frac{m+1}{2}) \; \Gamma(\frac{m}{2}+2)}{\Gamma(\frac{m}{2}+1) \; \Gamma(\frac{m+3}{2})} \quad \frac{(cE_o/BL)}{Dk_{max}^{\;2}} \tag{10}$$

For an inverse power law with spectral index $n = 2$ and a $\cos^2\theta$ angular ($m = 2$) dependence we find $k_o \propto L^{-1}$. This scaling of outer scale wavenumber k_o with the initial Pedersen conductivity gradient scale length when $n \simeq 2-3$ is consistent with recent numerical simulation studies of the $\underline{E} \times \underline{B}$ gradient drift instability in ionospheric plasma clouds [Keskinen et al., 1980a; Keskinen and Ossakow, 1980]. Furthermore, by taking $D = 1\ m^2/sec$, $L = 6$ km, $V_o = cE_o/B = 100$ m/sec, $2\pi/k_{max} = 20$ m [McDonald et al., 1978; Baker and Uiwick, 1978] we find from (10) that $2\pi/k_o \simeq 200$ m.

We further note that these results may also be applicable to the natural ionosphere since if one replaces the terms $D\nabla^2\delta\Sigma$ and $\underline{E}_o \cdot \nabla\delta\Sigma$ in (4) and (5) with $-\nu_R\delta\Sigma$ and $\left(g \times \hat{\underline{z}}/\nu_{in}\right) \cdot \nabla\delta\Sigma$,

respectively, where ν_R is a recombination rate, $\underline{g}$ denotes gravity, and ν_{in} the collision frequency with neutrals, then (4) and (5) are similar to the equations that describe the collisional Rayleigh-Taylor instability [Ossakow et al., 1979] in the equatorial F layer. The adaptation of this method to the collisional Rayleigh-Taylor instability and current convective instability is now in progress.

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